# Squeezed gluon vacuum and the global colour model of QCD

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## Abstract

We discuss how the vacuum model of Celenza and Shakin with a squeezed gluon condensate can explain the existence of an infrared singular gluon propagator frequently used in calculations within the global colour model. In particular, it reproduces a recently proposed QCD-motivated model where low energy chiral parameters were computed as a function of a dynamically generated gluon mass. We show how the strength of the confining interaction of this gluon propagator and the value of the physical gluon condensate may be connected.

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The Global Colour Model (CGM) of Quantum Chromodynamics (QCD), whose main aspects have been reviewed in the last years [1] - [3], is a quark-gluon quantum field theory that very successfully models QCD for low energy hadronic processes. In this approach an effective gluon correlator models the interaction between quark currents, and quark and gluon confinement may appear via the criterion of absence of real $q^2$ poles for the propagators [4,5]. This is, for instance, the case of an effective gluon propagator with an infrared singularity like a delta function $\delta(k)$ at low energy [5]. There are many recent calculations exemplifying the remarkable success of this procedure [6,7]. It relates the hadronic properties to the Schwinger functions of quarks and gluons, therefore, when comparing the theoretical calculations to some low energy data, as pseudoscalars masses and decay constants or other chiral parameters, we are learning how is the infrared behavior of the quark and gluon propagators. As the time goes on this semi phenomenological tool may reveal to be even more successful than the relativistic quark model or the bag model. However, the question for the mechanism which leads to the infrared singularities present in the gluon propagator in such calculations remains open.

The infrared enhancement of the gluon propagator due to the nonAbelian character of the theory and in particular due to the gluon-gluon self coupling, in principle could be understood in a rigorous study of the QCD vacuum. Unfortunately the simple perturbative vacuum is unstable [8], and there is no stable (gauge invariant) coherent vacuum in Minkowski space [9]. On the other hand in the context of the construction of a gauge invariant, stable QCD vacuum in Minkowski space, the squeezed condensate of gluons has become a topic of interest to uncover the underlying dynamics of the theory [10] - [12]. Within this class of vacuum models, Pavel *et al.* [13] recently proposed a phenomenological vacuum based on Abelian QCD which leads exactly to the infrared singularity in the gluon propagator as considered in Ref. [5]. On the other hand calculations within the GCM suggest that the best fit of chiral symmetry breaking parameters are obtained with a propagator containing a delta function plus a propagator that behaves as $1/k^2$ in the ultraviolet (consistent with QCD) but damped at $k^2 = 0$ [7,14]. Unfortunately, this is not the case of Ref. [13] where a simple $1/k^2$ is obtained together with the delta function.

Even if a model of the squeezed QCD vacuum is not determined from first principles, its properties may be very representative if they lead to a consistent phenomenology, indicating the path to the actual vacuum. In this note we show that the model of the QCD vacuum proposed by Celenza and Shakin [10,11] reproduces completely the gluon propagator of Ref. [14], *i.e.* it gives a singular part *à la* Munczek and Nemirovsky [5] as well as a piece containing a dynamically generated mass [15,16]. The effective dynamical gluon mass is the unique parameter in the model, and the gluon propagator obtained according to Ref. [10,11] is totally compatible with the idea of the GCM. We will briefly discuss the CGM and justify a gluon correlator which reproduces many aspects of the chiral symmetry breaking phenomenology. Then, we show that this gluon correlator naturally appears in the squeezed vacuum model of Ref. [10,11], verifying that the parameters involved in the model (the strength of the confining interaction, the dynamical gluon mass and the value of the physical gluon condensate) are consistent with the ones in the literature. The coincidence between the models may indeed suggest an interesting role for the Celenza and Shakin vacuum model.

The action of the GCM can be obtained from the QCD generating functional through the standard method presented in Ref. [1]:

$$Z[\bar{\eta},\eta] = N\int D\bar{q}DqDA exp\left(-S[\bar{q},q,A] + \int d^4x(\bar{\eta}q + \bar{q}\eta)\right), \tag{1}$$

where

$$S[\bar{q},q,A] = \int d^4x\left(\bar{q}(\not{\partial} + m - \imath g\frac{\lambda^a}{2}\not{A}^a)q + \frac{1}{4}G^a_{\mu\nu}G^a_{\mu\nu}\right), \tag{2}$$

and $G^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + gf^{abc}A^b_\mu A^c_\nu$. In the above equation we have not written the gauge fixing term, the ghost field term and its integration measure. Introducing a source term for the gauge field and writing

$$exp(W[J]) = \int DA exp\left(-\int d^4x(\frac{1}{4}G^a_{\mu\nu}G^a_{\mu\nu} - J^a_\mu A^a_\mu)\right), \tag{3}$$

the generating functional becomes

$$Z[\bar{\eta},\eta] = N\int D\bar{q}Dq exp(-\bar{q}(\not{\partial}+m)q+\bar{\eta}q+\bar{q}\eta) exp\left(W\left[\imath g\bar{q}\frac{\lambda^a}{2}\gamma_\mu q\right]\right), \qquad (4)$$

where the spacetime integration is implied. The functional $W[J]$ is the generator of connected gluon $n$-point functions without quark-loop contributions. It may be written as

$$W[J] = \frac{1}{2}\int d^4x d^4y J^a_\mu(x) g^2 D_{\mu\nu}(x-y) J^a_\nu(y) + W_R[J]. \qquad (5)$$

The main characteristic of the GCM is the fact that we neglect the higher $n$-point functions contained in Eq.(5) and expressed by $W_R[J]$. The effect of this approximation can only be measured in the model building, but it is expected that the phenomenological propagator $g^2D_{\mu\nu}(x-y)$ retain most of the information about the non-Abelian character of QCD. With this approximation the generating functional can be factorized as

$$Z[\bar{\eta},\eta] = exp\left(W_R\left[\imath g\frac{\delta}{\delta\eta}\frac{\lambda^a}{2}\gamma_\mu\frac{\delta}{\delta\bar{\eta}}\right]\right) Z_{GCM}[\bar{\eta},\eta]. \qquad (6)$$

$Z_{GCM}$ is giving by

$$Z_{GCM}[\bar{\eta},\eta] = N\int D\bar{q}Dq exp\left(-S_{GCM}[\bar{q},q]+\bar{\eta}q+\bar{q}\eta\right), \qquad (7)$$

with

$$S_{GCM}[\bar{q},q] = \int d^4x\bar{q}(x)(\not{\partial}+m)q(x) + \frac{1}{2}\int d^4x d^4y J^a_\mu(x) g^2 D_{\mu\nu}(x-y) J^a_\nu(y). \qquad (8)$$

The action $S_{GCM}[\bar{q},q]$ together with the generating functional $Z_{GCM}[\bar{\eta},\eta]$ defines the GCM. The idea of the model is that the nonperturbative behavior that could be missed in the truncation performed above can be mostly represented by an effective model in the infrared of the gluon propagator in Eq.(8).

There are several discussions in the literature about the nonperturbative behavior of the gluon propagator [2,17,16], as well as ansätze motivated by an impressive fitting of the low energy QCD phenomenology. One such case is the propagator proposed by Frank and Roberts [7] which yields the expected QCD behavior in the ultraviolet and presents an integrable singularity at the origin. This is accomplished in Landau gauge by the following form (in the sequence all the momenta are in Euclidean space)

$$g^2 D_{\mu\nu}(k) = \left\{ \delta_{\mu\nu} - \frac{k_\mu k_\nu}{k^2} \right\} \frac{g^2}{k^2[1+\Pi(k^2)]} \ , \tag{9}$$

where,

$$\begin{aligned} \Delta(k^2) &\equiv \frac{g^2}{k^2[1+\Pi(k^2)]} \\ &= 4\pi^2 d \left[ 4\pi^2 m_t^2 \delta^4(k) + \frac{1-e^{(-k^2/[4m_t^2])}}{k^2} \right] \ , \end{aligned} \tag{10}$$

with $d = 12/(33-2n_f)$, and $n_f = 3$ (considering only three quark flavors). The mass scale $m_t$ determined in Ref. [7] was interpreted as marking the transition between the perturbative and nonperturbative domains. We recently modified the above gluon propagator in order to corroborate with the suggestion of several works [15,16] that the gluon can develop a dynamical mass and established that the unique parameter, $m_t$, could conveniently be substituted by the scale associated to this mass. Then, the Eq.(10) should be modified to [14]

$$\begin{aligned} \Delta(k^2) &\equiv \frac{g^2}{k^2[1+\Pi(k^2)]} \\ &= 4\pi^2 d \left[ 4\pi^2 m_g^2 \delta^4(k) + \frac{1}{k^2 + m^2(k^2)} \right], \end{aligned} \tag{11}$$

where,

$$m^2(k^2) = m_g^2 \frac{m_g^2}{k^2 + m_g^2} \ , \tag{12}$$

is the dynamical gluon mass. Eq.(12) interpolates between the large momenta behavior of the gluon mass predicted by the operator product expansion [18]

$$m_g^2(k^2) \sim \frac{34N\pi^2}{9(N^2-1)} \frac{\left\langle \frac{\alpha_s}{\pi} G^{\mu\nu} G_{\mu\nu} \right\rangle}{k^2}, \tag{13}$$

and its finite constant infrared behavior as discussed in Ref. [15]. In Eq.(13) $\left\langle \frac{\alpha_s}{\pi} G^{\mu\nu} G_{\mu\nu} \right\rangle$ is the gluon condensate. Eq.(11) has in its first term the singularity of Eq.(10), and the dynamical gluon mass gives a natural (and needed) damping of the second term at small $k^2$. In Ref. [14] we obtained a good fit to some chiral parameters ( as the pion mass and the quark condensate) using the value $m_g \sim 600 \ MeV$. Our next step is to show how the model of Ref. [10,11] can fully describe the behavior of Eq.(11).

Colour-singlet coherent states can be formed integrating over the group elements associated with the colour group

$$|\mathbf{Z}_o(t) >= N \int [dg] U(g) |\mathbf{Z}(t) >, \tag{14}$$

where $N$ is a numerical factor, $g$ specify an element of the group $SU(N)$ and $U(g)$ are rotation operators whose properties are described in Ref. [11]. Celenza *et al.* [11] were able to form a squeezed state $|\mathbf{Z}_o >$ such that

$$< \mathbf{Z}_o| : \mathbf{E}_a(\mathbf{r},t) \cdot \mathbf{E}_a(\mathbf{r},t) : |\mathbf{Z}_o >= 0, \tag{15}$$

and

$$\frac{g^2}{2} < \mathbf{Z}_o| : \mathbf{B}_a(\mathbf{r},t) \cdot \mathbf{B}_a(\mathbf{r},t) : |\mathbf{Z}_o > \neq 0. \tag{16}$$

Meaning that the condensate is purely magnetic ($\mathbf{B}_a(\mathbf{r},t)$), while the color-electric field ($\mathbf{E}_a(\mathbf{r},t)$) leads to a vanishing condensate.

Many of the properties of this model were obtained previously with the assumption that the gluon field could be decomposed into a constant condensate field $\mathcal{G}^a_\mu$ and the quantum fluctuations $g^a_\mu(x)$ around it [10],

$$G^a_\mu(x) = \mathcal{G}^a_\mu + g^a_\mu(x). \tag{17}$$

Taking into account Eq.(17) and $< g^a_\mu(x) >= 0$, one obtains the decomposition of the nonperturbative gluon propagator into two parts,

$$g^2 G^{ab}_{\mu\nu}(x-y) =< g^2 \mathcal{G}^a_\mu \mathcal{G}^b_\nu > + < g^2 g^a_\mu(x) g^b_\nu(y) > . \tag{18}$$

Note that the choice of Eq.(17) is a purely phenomenological one, and the correct construction of the model goes through the steps of Ref. [11]. However, in this way it is easier to verify some of the consequences of the model. First, it does generate a massive effective lagrangian [10]

$$\mathcal{L}_{m_g}(x) = -\frac{1}{4} G^a_{\mu\nu}(x) G^{\mu\nu}_b(x) + \frac{m_g^2}{2} g^\mu_a(x) g^a_\mu(x) + ... \tag{19}$$

Secondly, the gluon mass $m_g^2$ is equal to $(15/32) < g^2\mathcal{G}^2 >$. In Ref. [10] the square of this condensate has been identified with the phenomenological gluon condensate of Shifman *et al.* [19], but we do not pursue this point since the gluon mass may contain a large part of dynamics, and we do not see any *a priori* reason for the identity of the condensates (although it looks phenomenologically plausible). Finally, what more interest us is that in Euclidean momentum space the effective nonperturbative gluon propagator corresponding to the decomposition of Eq.(18) has the form [20]

$$g^2\Delta(k^2) \propto \left[16\pi^4 m_g^2 \delta^4(k) + \frac{1}{k^2+m_g^2}\right], \tag{20}$$

where the $\delta^4(k)$ comes from the condensate field and the massive propagator is a consequence of both fields (since the condensate field generates the gluon mass). Once we have not saturated the condensate resulting from Eq.(18) with the one of Shifman *et al.*, the value of the dynamical gluon mass can be considered a free parameter and not yet related to the gluon condensate. Note that we do have some freedom in the definition of the factors of Eq.(11) [14], as well as in the definition of the gluon mass in Eq.(20), this is why we may affirm that Eq.(11) and Eq.(20) are totally compatible. The important point is that the form of the propagator in Eq.(20) is consistent with the propagator of Ref. [14], and gives a meaning for the confining propagators of the GCM. It is interesting that the proposal of Ref. [14] was to modify the propagator of Ref. [7] introducing the concept of the dynamical gluon mass, and it turned out that this modification just led to the model of Ref. [10].

In the Celenza and Shakin model the gluon mass and the strength of the confining interaction are related to the constant condensate field, and the square of this one was identified with the gluon condensate of Shifman *et al.* [19]. As we said above it is not clear to us how much of the dynamics can be built over this condensate, and we can let the gluon mass as a free parameter in Eq.(20). However, we would like to show that this identity is natural. There are several ways to verify that the gluon mass (or the condensate field) are related to the phenomenological gluon condensate. We obviously expect that the dynamical gluon mass is connected to the gluon condensate through the operator product expansion as described by Eq.(13). We can also compute the vacuum energy as a function of the propagator in Eq.(20), and than equalize this vacuum energy to the trace anomaly expression $\frac{\beta(g)}{2g}\langle G_{\mu\nu}G^{\mu\nu}\rangle$, where $\beta(g)$ is the perturbative $\beta$ function of the renormalization group equation. This approach was used in the last paper of Ref. [15] and more recently in Ref. [17], where the relation with the gluon condensate was explored to fix the gluon mass. As these methods have already been discussed for massive gluons, we propose here other procedures where we can show that the gluon propagator of Eq.(20) can indeed be related to the gluon condensate as well as to other typical hadronic parameters. We start by showing the relation between the confining interaction and the bag constant.

The bag constant ($\mathcal{B}$) for a given gluon propagator can be obtained through [21]

$$\mathcal{B} = 12\pi^2 \int \frac{k^2 dk^2}{(2\pi)^4} \left( \ln \left( \frac{A^2(k^2)k^2 + B^2(k^2)}{A^2(k^2)k^2} \right) - \frac{B^2(k^2)}{A^2(k^2)k^2 + B^2(k^2)} \right), \tag{21}$$

where $A(k^2)$ and $B(k^2)$ appear in the inverse of the renormalized quark propagator

$$S^{-1}(k) = \imath \not{k} + \Sigma(k) = \imath \not{k} A(k^2) + B(k^2). \tag{22}$$

The form of $A(k^2)$ and $B(k^2)$ is obtained solving the Schwinger-Dyson equations for a given gluon propagator, and here enter the information about the GCM propagators.

The Schwinger-Dyson equations for a gluon propagator equal to Eq.(11) were discussed in Ref. [14]. To avoid a numerical calculation that may conceal the simplicity of the problem, we will approximate the gluon propagator only by its confining part. The part proportional to $(k^2 + m_g^2)^{-1}$ is responsible for the tail of the propagator, actually, the larger is the gluon mass the smaller is the contribution to this tail. Therefore, we will not introduce a large error in neglecting the second term of Eq.(11) or Eq.(20). Proceeding in this way it is easy to verify that the propagator $\Delta(k^2) = 16\pi^4 m_g^2 \delta^4(k)$ imply in the following solution for $A(k^2)$ and $B(k^2)$ when $k^2 < m_g^2$ [2]

$$A(k^2) = 2 \quad , \quad B(k^2) = 2\sqrt{m_g^2 - k^2}. \tag{23}$$

For $k^2 > m_g^2$ we have $B(k^2) = 0$. With the solution of Eq.(23) the calculation of $\mathcal{B}$ is straightforward and gives

$$\mathcal{B} = \frac{m_g^4}{16\pi^2}. \tag{24}$$

For a gluon mass of approximately $600\,MeV$ [14,15] we obtain $\mathcal{B} = (169\,MeV)^4$ which, considering the approximation performed to obtain Eq.(23), is in good agreement with the MIT value of $(146\,MeV)^4$.

The bag constant can be related to the string tension for fermions in the fundamental representation through [22]

$$K_F = (8\pi\alpha_s C_F \mathcal{B})^{1/2}, \tag{25}$$

where $C_F$ is the quadratic Casimir operator for the fundamental representation. On the other hand this same string tension has been estimated in the last paper of Ref. [15] as

$$K_F \approx \frac{\pi^3}{9} \frac{\left\langle \frac{\alpha_s}{\pi} G^{\mu\nu} G_{\mu\nu} \right\rangle}{m_g^2}. \tag{26}$$

From the identity of Eq.(25) and Eq.(26) we obtain

$$m_g^4 = \frac{\pi^4}{9} \left( \frac{3}{2\pi\alpha_s} \right)^{1/2} \left\langle \frac{\alpha_s}{\pi} G^{\mu\nu} G_{\mu\nu} \right\rangle. \tag{27}$$

Eq.(27) shows the connection between the gluon mass and the gluon condensate. Taking $\alpha_s \approx 1$ and $\left\langle \frac{\alpha_s}{\pi} G^{\mu\nu} G_{\mu\nu} \right\rangle \simeq (0.01)\ GeV^4$ [19] we obtain $m_g \approx 550\ MeV$, which is consistent with the value we quoted before (see Refs. [14,10]).

Another way to verify the relation between the gluon mass and the gluon condensate is comparing the bag energy density to minus the vacuum expectation value of the trace of the QCD energy momentum tensor, although it is known that this comparison does not give reasonable numbers [23]. The energy density in the bag model at the equilibrium radius is equal to

$$\Omega = 4\mathcal{B}, \tag{28}$$

and from the vacuum expectation value of the trace of the QCD energy momentum tensor we obtain

$$\Omega = \frac{1}{4}\left\langle \Theta_{\mu\mu} \right\rangle = -\frac{1}{24}\frac{\alpha_s}{\pi}(11N - 2n_f)\left\langle G_{\mu\nu}G^{\mu\nu} \right\rangle . \tag{29}$$

which gives

$$m_g^4 = \frac{(11N - 2n_f)\pi^2}{6}\left\langle \frac{\alpha_s}{\pi} G^{\mu\nu}G_{\mu\nu} \right\rangle . \tag{30}$$

As expected we obtain a gluon mass proportional to the gluon condensate, but about 1.5 times the desired value (considering $n_f = 3$).

In conclusion, we discussed how the QCD vacuum model of Celenza and Shakin [11] explain the presence of a confining gluon propagator frequently used in calculations within the global colour model. In particular, it reproduces the propagator of Ref. [14] where the dynamically generated gluon mass plays a fundamental role. We considered the strength of the confining interaction, equal to the gluon mass, to be a free parameter contrarily to Ref. [10], based on the fact that the dynamics may modify the gluon propagator. We have given arguments showing that the gluon mass is indeed related to the phenomenological gluon condensate, as well as it can be related to the string tension and bag constant. The coincidence between the propagators used in applications of the global colour model and the one resulting from the squeezed vacuum model of Celenza and Shakin might indicate some intrinsic property of the actual QCD vacuum.

## ACKNOWLEDGMENTS

I am grateful to A. E. Dorokhov for valuable remarks. This research was supported in part by the Conselho Nacional de Desenvolvimento Cientifico e Tecnológico (CNPq), Fundação de Amparo a Pesquisa do Estado de São Paulo (FAPESP), and by Programa de Apoio a Núcleos de Excelência (PRONEX).

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